RELEVANCE OF HIGH-MOMENTUM NUCLEONS FOR NUCLEAR PHENOMENA

W. H. Dickhoff,[1, 2, †] E. P. Roth,[2] and M. Radici[3]

[1] *Laboratory of Theoretical Physics, University of Gent, Proeftuinstraat 86, B-9000 Gent, Belgium*

[2] *Department of Physics, Washington University, St. Louis, Missouri 63130, USA*

[3] *Istituto Nazionale di Fisica Nucleare, Sezione di Pavia, Pavia, Italy, I-27100*

A brief review is given concerning the status of the theoretical work on nucleon spectral functions. A recent concern about the validity of the concept of spectroscopic factors as deduced from (e,e′p) reactions at higher Q^2, is discussed in some detail. The consequences of the observed spectral strength are then considered in the context of nuclear saturation. It is argued that short-range correlations are mainly responsible for the actual value of the observed charge density in ^{208}Pb and by extension for the empirical value of the saturation density of nuclear matter. This observation combined with the general understanding of the spectroscopic strength suggests that a renewed study of nuclear matter, emphasizing the self-consistent determination of the spectral strength due to short-range and tensor correlations, may shed light on the perennial nuclear saturation problem. First results using such a scheme are presented.

1. INTRODUCTION

During the last fifteen years considerable progress has been made in elucidating the limits of the nuclear mean-field picture. The primary tool in exhibiting these limits in a quantitative fashion has been provided by the (e,e′p) reaction [1, 2, 3, 4]. In this paper the status of the theoretical understanding of the spectroscopic factors that have been deduced from the analysis of this reaction will be reviewed briefly. The qualitative features of the strength distribution can be understood by realizing that a considerable mixing occurs between hole states and two-hole one-particle (2h1p) states. This leads to the observed fragmentation pattern which exhibits a single peak for valence hole states near the Fermi energy, albeit with reduced strength. A strongly fragmented strength distribution is observed for more deeply bound states. This fragmentation is due to the strong coupling to 2h1p states and the presence of these states at energies corresponding to these more deeply bound hole states which leads to correspondingly small energy denominators. For quantitative results one also requires the inclusion of short-range and tensor correlations. On the one hand, this leads to a global depletion of mean-field orbitals which ranges from 10% in light nuclei to about 15% in heavy nuclei and nuclear matter [5, 6]. This depletion effect, on the other hand, is then compensated by the admixture of high-momentum components in the ground state. These high-momentum nucleons have not yet been unambiguously identified experimentally using the (e,e′p) reaction. The search for these high-momentum components in valence states has not been successful [7, 8], as was anticipated by earlier theoretical work [9].

A recent publication [10] has challenged the conventional interpretation of the (e,e′p) reaction with regard to valence hole states. This challenge consists in questioning the validity of the constancy of the spectroscopic factor as a function of the four-momentum, Q^2, transferred by the virtual photon to the knocked-out nucleon. This recent work employs a Skyrme-Hartree-Fock bound-state wave functions for the initial proton, a Glauber-type description of the final-state interaction of the outgoing proton, and a factorization approximation for the electromagnetic vertex for the description of the reaction at higher Q^2. The results obtained in Ref. [10] display an increasing spectroscopic strength with increasing Q^2 for the ^{12}C nucleus. While the theoretical definition of the spectroscopic factor is unambiguously independent of the probe, it is worth studying the description of the data at higher Q^2 in a consistent manner. To this end we employ a recently developed eikonal approximation [11, 12, 13, 14, 15] to describe the outgoing proton under conditions

†Electronic address: wimd@wuphys.wustl.edu; URL: `http://www.physics.wustl.edu/~wimd`

appropriate for a recent JLab experiment [16]. This description of the final-state-interaction (FSI) is combined with previous results for the quasihole wave functions obtained for ^{16}O [17] which were employed for the description [18] of a low Q^2 experiment [19]. The absorption of the outgoing proton in the eikonal approach is then related to the corresponding absorption experienced by a nucleon in nuclear matter as obtained from the self-energy. This self-energy is obtained from self-consistent calculations of the nucleon spectral functions including the effect of short-range and tensor correlations in nuclear matter [20]. This approach for the analysis at $Q^2 = 0.8(\mathrm{GeV/c})^2$ is discussed in this contribution and some initial results are presented.

Based on the experimental and theoretical results for the spectral strength distribution, one may wonder what the consequences are for the "energy" or "Koltun" sum rule [21, 22, 23]. In principle, one can ascertain that a perfect agreement of the theoretical strength with the experimental one, will yield a correspondingly good agreement for the energy per particle provided three-body forces are not too important. Initial indications of the relevance of high-momentum nucleons, which sofar have not been observed directly, for the question of the energy per particle have already been raised in earlier work [17]. In addition, we will argue in this work that the actual value of the nuclear sturation density is dominated by the effects of short-range and tensor correlations (SRC). Recent experimental work will be discussed which supports this claim. Based on these considerations, it is sugested that a renewed study of the nuclear satuation problem is in order. Results will be discussed which suggests that new insights may be obtained using this approach. We close with some conclusions.

2. STATUS OF THEORETICAL RESULTS FOR SPECTROSCOPIC STRENGTH

One of the critical experimental ingredients in clarifying the nature of nuclear correlations has only become available over the last decade and a half. It is therefore not surprising that all schemes that have been developed to calculate nuclear matter saturation properties are not based on the insights that these experiments provide. Before discussing the implications of these insights, we will review these results in this section. Exclusive experiments, involving the removal of a proton from the nucleus which is induced by a high-energy electron that is detected in coincidence with the removed proton, have given access to absolute spectroscopic factors associated with quasihole states for a wide range of nuclei. [1, 2, 3, 4] The experimental results indicate that the removal of single-particle strength for quasihole states near the Fermi energy corresponds to about 65%. The spectroscopic factors obtained in these experiments can be directly related to the single-particle Green's function of the system which is given by

$$
\begin{aligned}
g(\alpha,\beta;\omega) = & \sum_m \frac{\left\langle \Psi_0^A \middle| a_\alpha \middle| \Psi_m^{A+1} \right\rangle \left\langle \Psi_m^{A+1} \middle| a_\beta^\dagger \middle| \Psi_0^A \right\rangle}{\omega - (E_m^{A+1} - E_0^A) + i\eta} \\
& + \sum_n \frac{\left\langle \Psi_0^A \middle| a_\beta^\dagger \middle| \Psi_n^{A-1} \right\rangle \left\langle \Psi_n^{A-1} \middle| a_\alpha \middle| \Psi_0^A \right\rangle}{\omega - (E_0^A - E_n^{A-1}) - i\eta}.
\end{aligned} \tag{1}
$$

This representation of the Green's function is referred to as the Lehmann-representation and involves the exact eigenstates and corresponding energies of the A- and $A \pm 1$-particle systems. Both the addition and removal amplitude for a particle from (to) the ground state of the system with A particles must be considered in Eq. (1). Only the removal amplitude has direct relevance for the analysis of the (e,e′p) experiments. The spectroscopic factor for the removal of a particle in the single-particle orbit α, while leaving the remaining nucleus in state n, is then given by

$$
z_\alpha = \left| \left\langle \Psi_n^{A-1} \middle| a_\alpha \middle| \Psi_0^A \right\rangle \right|^2, \tag{2}
$$

which corresponds to the contribution to the numerator of the second sum in Eq. (1) of state n for the case $\beta = \alpha$. Another important quantity, which also contains this information, is the spectral function associated with single-particle orbit α. The part corresponding to the removal of particles, or hole spectral function, is given by

$$
S_h(\alpha,\omega) = \sum_n \left| \left\langle \Psi_n^{A-1} \middle| a_\alpha \middle| \Psi_0^A \right\rangle \right|^2 \delta(\omega - (E_0^A - E_n^{A-1})), \tag{3}
$$

which corresponds to the imaginary part of the diagonal elements of the propagator and characterizes the strength distribution of the single-particle state α as a function of energy in the $A-1$-particle system. From this quantity one can therefore obtain another key ingredient that gauges the effect of correlations, namely the occupation number which is given by

$$n(\alpha) = \int_{-\infty}^{\epsilon_F} d\omega \; S_h(\alpha,\omega) = \left\langle \Psi_0^A \middle| a_\alpha^\dagger a_\alpha \middle| \Psi_0^A \right\rangle . \tag{4}$$

In the experimental analysis the quantum number α is related to the actual Woods-Saxon potential required to both reproduce the correct energy of the hole state as well as the shape of the corresponding (e,e′p) cross section for this particular transition. The remaining parameter required to fit the actual data then becomes the spectroscopic factor associated with this transistion. In this analysis the reduction of the flux associated with the scattering of the outgoing proton is incorporated by the use of empirical optical potentials describing elastic proton-nucleus scattering data. Experiments on ^{208}Pb result in a spectroscopic factor of 0.65 for the removal of the last $3s_{1/2}$ proton [2]. Additional information about the occupation number of this orbit can be obtained by analyzing elastic electron scattering cross sections of neighboring nuclei [24]. The actual occupation number for the $3s_{1/2}$ proton orbit obtained from this analysis is about 10% larger than the quasihole spectroscopic factor [25, 26]. A recent analysis of the (e,e′p) reaction on ^{208}Pb in a wide range of missing energies and for missing momenta below 270 MeV/c yields information on the occupation numbers of more deeply bound orbitals. The data suggest that all deeply bound orbits are depleted by the same amount of about 15% [27, 28].

As discussed in the introduction, the general properties of the experimental strength distributions can be understood on the basis of the coupling between single-hole states and 2h1p states. This implies that a proper inclusion of this coupling in the low-energy domain is required in theoretical calculations that aim at reproducing the experimental distribution of the strength. Such calculations have been successfully performed for medium-heavy nuclei [29, 30]. Indeed, calculations for the strength distribution for the removal of protons from ^{48}Ca demonstrate that an excellent qualitative agreement with the experimental results is obtained when the coupling of the single-hole states to low-lying collective states is taken into account [30]. This coupling is taken into account by calculating the microscopic RPA phonons and then constructing the corresponding self-energy. The solution of the Dyson equation then provides the theoretical strength distribution. By adding the additional depletion due to short-range correlations a quantitative agreement is obtained although no explicit calculation for these nuclei including both effects has been performed to date. The corresponding occupation numbers calculated for this nucleus also indicate that the influence of collective low-lying states, associated with long-range correlations, on the occupation numbers is confined to single-particle states in the immediate vicinity of the Fermi level. The description of the spectroscopic strength in ^{16}O is not as successful [31] on account of the complexity of the low-energy structure of this nucleus. Attempts to describe the proper inclusion of microscopic particle-particle and particle-hole phonons in a Faddeev approach for this nucleus are currently in progress [32, 33].

For a global understanding of the strength distribution it is also necessary to account for the appearance of single-particle strength at high momenta as a direct reflection of the influence of short-range correlations. These high-momentum nucleons make up for the missing strength that has been documented in (e,e′p) experiments. Results for ^{16}O [9, 17] corroborate the expected occupation of high-momenta but put their presence at high missing energy. This can be understood in terms of the admixture of a high-momentum nucleon requiring 2h1p states which can accomodate this momentum through momentum conservation. Since two-hole states combine to small total pair momenta, one necessarily needs a high-momentum nucleon (of about equal and opposite value to the component to be admixed) with corresponding high excitation energy. As a result, one expects to find high-momentum components predominantly at high missing energy. Recent experiments at JLab attempt at a quantitative assessment of the strength of these high-momentum nucleons [34].

3. CONSISTENT APPROACH TO THE ANALYSIS OF SPECTROSCOPIC STRENGTH AT HIGH Q^2

Recent work [10] has challenged the conventional interpretation of the (e,e′p) reaction with regard to valence hole states. An analysis of all data for the ^{12}C nucleus at low Q^2 is obtained in this work using the standard analysis technique developed by the NIKHEF group. For experiments at higher Q^2 a different approach was used. For example, Skyrme-Hartree-Fock bound-state wave functions are employed for which no spectroscopic factor at NIKHEF kinematics is quoted. To account for final-state interactions of the proton, which absorbs the virtual photon, a Glauber-type description is employed with some slight adjustments of the input nucleon-nucleon cross section to account for in-medium effects [35]. As is usual in this approach, a factorization approximation for the electromagnetic vertex is employed. It is unclear whether this represents a serious approximation. At low Q^2, it would be unacceptable. The results of Ref. [10] exhibit an increase in the spectroscopic strength with increasing Q^2 for the ^{12}C nucleus.

The theoretical definition of the spectroscopic factor (see Eq. (2)) involves a matrix element of a particle removal operator between the ground state and an appropriate state in the system with one particle less. Clearly, no dependence on Q^2 can be generated by such an expression. The extraction of spectroscopic factors clearly involves a detailed model describing the exit of a strongly interacting particle from the nucleus. It is therefore of great interest to pursue the question whether one can extend the analysis involving spectroscopic factors to kinematical conditions involving higher Q^2. We have therefore developed an approach to test this possibility which relies on a recently developed eikonal model of the FSI [11, 12, 13, 14, 15]. This approach will be used to describe a recent JLab experiment on ^{16}O [16] at $Q^2 = 0.8$ (GeV/c)^2. We proceed by describing some of the key features of this approach before discussing the other ingredients of this calculation.

The nuclear response in exclusive (e,e′p) process can be parametrized as a bilinear product of matrix elements of the different helicity components of the nuclear current, which describe the transition from the initial to the final hadronic states. In the projection-operator approach and within the framework of the Distorted-Wave Impulse Approximation (DWIA) [36], it is possible to project out of the total Hilbert space a suitable channel where the matrix elements are written in a one-body representation as

$$J^{\lambda}_{\alpha s' s}(\vec{p}_f, \vec{q}) = \int d\vec{r} d\sigma e^{\mathrm{i}\vec{q}\cdot\vec{r}} \chi^{(-)*}_{\vec{p}_f s'}(\vec{r},\sigma) \hat{J}_{\lambda}(\vec{q},\vec{r},\sigma) \phi_{\alpha s}(\vec{r},\sigma) \,, \tag{5}$$

where $\vec{q}$ is the momentum carried by the virtual photon and $\vec{p}_f, s'$ are the momentum and spin of the detected nucleon, leaving a hole in the residual nucleus with collective quantum numbers α and spin s. If the detected nucleon is moving fast, the scattering wave function χ, with incoming-wave boundary conditions, is usually approximated by an eikonal wave through the Glauber method [37]. However, since for a fastly moving object the nuclear density can be considered roughly constant inside all the nuclear volume (except for a small portion on the surface), the eikonal wave function can be further approximated in lowest order by a plane wave with complex momentum $\vec{P}_f = \vec{p}_f + \mathrm{i}\vec{p}_I$ and $\vec{p}_I \parallel \vec{p}_f$:

$$J^{\lambda}_{\alpha s' s}(\vec{p}_f, \vec{q}) = \sum_{\tilde{s}} \int d\vec{r} e^{\mathrm{i}\vec{q}\cdot\vec{r}} \left(e^{\mathrm{i}\vec{p}_f\cdot\vec{r}} e^{-\vec{p}_I\cdot\vec{r}} \right)^* \delta_{s'\tilde{s}} \langle \tilde{s} | \hat{J}_{\lambda}(\vec{q},\vec{r}) | s \rangle \phi_{\alpha s}(\vec{r}) \,. \tag{6}$$

The scattering wave function now corresponds to a uniformly damped plane wave, where the damping is driven by $\mathrm{Im}(\vec{P}_f) \equiv \vec{p}_I$. This corresponds to solving the Schrödinger equation with a complex potential for a particle travelling through homogeneous nuclear matter, i.e.

$$\left(\frac{-\hbar^2}{2m} \nabla^2 + \hat{V} + \mathrm{i}\hat{W} \right) \chi = E\chi \,, \tag{7}$$

or, equivalently,

$$\left(E - \hat{V} - \mathrm{i}\hat{W} \right) \chi = \frac{\vec{P}_f \cdot \vec{P}_f}{2m} \chi \; = \; \left(\frac{\vec{p}_f^{\,2} - \vec{p}_I^{\,2}}{2m} + \mathrm{i}\, \frac{\vec{p}_f \cdot \vec{p}_I}{m} \right) \chi \,, \tag{8}$$

from which a natural relationship between p_I and the absorptive part W of the potential is deduced. If the outgoing proton is sufficiently energetic, i.e. $p_f \geq 1$ GeV/c, and comes from a bound state with a momentum below the Fermi surface, this approximation has been shown to give reliable results [11, 12, 13] with a constant $p_I \propto W/p_f$.

The technical advantage of this approximation is that the representation of Eq. (6) in momentum space becomes completely analytical, provided that the integral is extended in the complex plane:

$$
\begin{aligned}
J^{\lambda}_{\alpha s' s}(\vec{p}_f, \vec{p}_I, \vec{q}) &= \sum_{\tilde{s}} \int d\vec{P} \delta(\vec{P}_f - \vec{P} - \vec{q}) \delta_{s'\tilde{s}} \langle \tilde{s} | \hat{J}_{\lambda}(\vec{q}, \vec{P}) | s \rangle \phi_{\alpha s}(\vec{P}) \\
&= \langle s' | \hat{J}_{\lambda}(\vec{q}, \vec{P}_f - \vec{q}) | s \rangle \phi_{\alpha s}(\vec{P}_f - \vec{q}) \, . \qquad (9)
\end{aligned}
$$

The conditions for an analytical extension are two: a suitable extension of the definition of the δ distribution in terms of complex variables, which automatically connects to the damped plane wave $e^{\mathrm{i}\vec{P}\cdot\vec{r}}$ as in the usual case with real variables; the integrand must be analytical and vanish asymptotically for $|\vec{P}| \to \infty$. Both conditions have been explored and verified in Ref. [14, 15].

The results presented for the E89003 kinematics have been obtained replacing $\vec{p}_I$ by the value predicted from nuclear-matter calculations (discussed in the next section) at the considered Q^2. This procedure corresponds to considering the wave function of the proton at the momentum corresponding to the kinematical setting. This wave function exhibits a damping due to the imaginary part of the self-energy. The dominant feature descibing this damping is obtained by associating a complex pole to the propagator which leads to the exponential damping of the wave function. This procedure can be used to obtain the complex part of the momentum [38] appropriate for this problem.

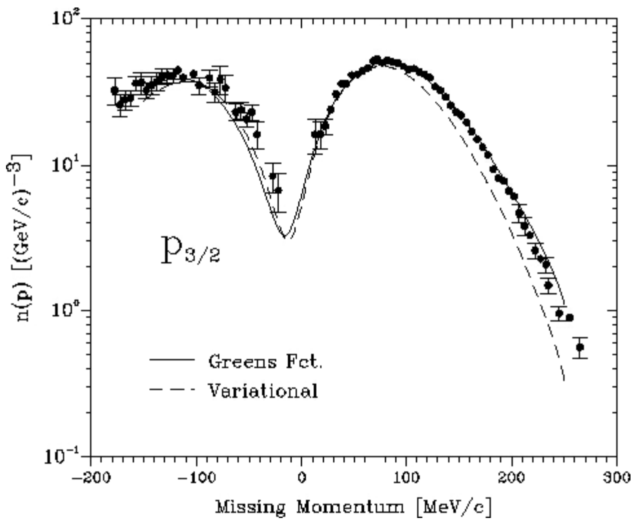

FIG. 1: Reduced cross section for the ^{16}O(e,e′p)reaction in parallel kinematics leading to the state at -6.32 MeV of the residual nucleus ^{15}N. Results of Green's function calculations [17] (solid line) are compared to those in the variational calculation of Ref. [39] (dashed line) and the experimental data [19].

We then start this analysis by employing results for the quasihole wave functions corresponding to $p_{3/2}$ (and $p_{1/2}$) protons from ^{16}O obtained in Ref. [17]. These wave functions have been used to extract spectroscopic factors [18] corresponding to a low Q^2 experiment performed at NIKHEF [19]. The result of this analysis is shown in Fig. 1. The description of the data using the Green's function calculation of the quasihole wave function require a spectroscopic factor of 0.537. With this reduction factor a good description of the data is obtained. Assuming the model for FSI described above and using the input from the nucleon self-energy in nuclear matter to describe the damping of the outgoing nucleon wave, we are in a position to calculate the cross section for the JLab experiment at $Q^2 = 0.8$ (GeV/c)^2 while keeping the **same** spectroscopic factors for the removal of the p-nucleons that generate a good description of the data at low Q^2 as shown in Fig. 1. The results of this calculation are shown in Fig. 2. These results suggest that it is indeed possible to use spectroscopic information obtained in the kinematical domain used at facilities like

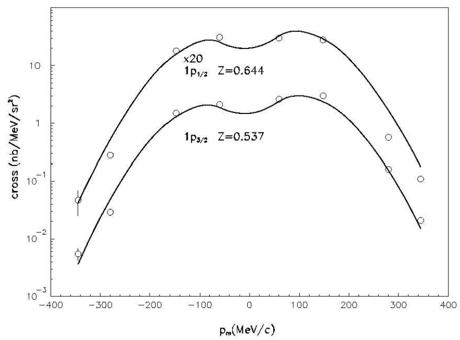

FIG. 2: Comparison of data obtained from the ^{16}O(e,e′p) reaction at $Q^2 = 0.8(\mathrm{GeV/c})^2$ [16] with a consistent theoretical approach involving the same spectroscopic factors as obtained for the NIKHEF experiment [19].

NIKHEF and, subsequently, correctly predict cross sections for the (e,e′p) reaction at higher Q^2. We emphasize that no adjustment of any input of the calculation has been performed to obtain the results in Fig. 2. Apparently, a consistent theoretical approach does not require a different interpretation of the spectroscopic strength under different kinematical conditions.

4. SELF-CONSISTENTLY DRESSED NUCLEONS IN NUCLEAR MATTER

Having confirmed the reliability of the interpretation of the spectroscopic strength obtained from the (e,e′p) reaction, it is possible to return to the discussion of the consequences of these results. We first collect some relevant information which will be used to argue that the empirical saturation density of nuclear matter is dominated by SRC. As discussed earlier, a recent analysis of the (e,e′p) reaction on ^{208}Pb up to 100 MeV missing energy and 270 MeV/c missing momenta indicates that all deeply bound orbits are depleted by the same amount of about 15% [27, 28]. This global depletion of the single-particle strength in about the same amount for all deeply bound states, as observed for ^{208}Pb, was anticipated [5, 40] on the basis of the experience that has been obtained with calculating occupation numbers in nuclear matter with the inclusion of SRC [41]. Such calculations suggest that about 15% of the single-particle strength in heavy nuclei is removed from the Fermi sea leading to the occupation of high-momentum states. This global depletion of mean-field orbitals can be interpreted as a clear signature of the influence of SRC. In turn, these results reflect on one of the key quantities determining nuclear saturation empirically. Elastic electron scattering from ^{208}Pb [42] clearly pinpoints the value of the central charge density in this nucleus. By multiplying this number by A/Z one obtains the relevant central density of heavy nuclei, corresponding to 0.16 nucleons/fm^3 or $k_F = 1.33$ fm^{-1}. Since the presence of nucleons at the center of a heavy nucleus is confined to s nucleons, and their depletion is dominated by SRC, one may conclude that the actual value of the saturation density of nuclear matter must also be closely linked to the effects of SRC. While this argument is particularly appropriate for the deeply bound $1s_{1/2}$ and $2s_{1/2}$ protons, it continues to hold for the $3s_{1/2}$ protons which are depleted predominantly by short-range effects (up to 15%) and by at most 10% due to long-range correlations as discussed above.

The binding energy of nuclei or nuclear matter usually includes only mean-field contributions to the kinetic energy when the calculations are based on perturbative schemes like the hole-line expansion. With the presence of high-momentum components in the ground state it becomes relevant to ask what the real kinetic and potential energy of the system look like in terms of the single-particle strength distributions. This theoretical result [21, 22] has the general form (Koltun sum rule)

$$E_0^A = \left\langle \Psi_0^A \right| \hat{H} \left| \Psi_0^A \right\rangle = \frac{1}{2} \sum_{\alpha\beta} \left\langle \alpha \right| T \left| \beta \right\rangle n_{\alpha\beta} + \frac{1}{2} \sum_{\alpha} \int_{-\infty}^{\epsilon_F} d\omega \; \omega S_h(\alpha, \omega) \tag{10}$$

in the case when only two-body interactions are involved. In this equation, $n_{\alpha\beta}$ is the one-body density matrix element which can be directly obtained from the single-particle propagator. A delicate balance exists between the repulsive kinetic energy term and the attractive contribution of the second term in Eq. (10) which samples the single-particle strength weighted by the energy parameter ω. When realistic spectral distributions are used to calculate these quantities surprising results emerge [17]. Such calculations for ^{16}O indicate that the contribution of the quasihole states to Eq. (10), corresponding to the $1s_{1/2}$, $1p_{3/2}$, and $1p_{1/2}$ orbitals, comprise only 37% of the total energy leaving 63% for the continuum terms that represent the spectral strength associated with the coupling to low-energy 2h1p states. These contributions therefore contain the presence of high-momentum components in the nuclear ground state reflecting the effect of SRC. Although these high momenta account for only 10% of the particles in the case of ^{16}O, their contribution to the energy is extremely important. These results give a first indication of the importance of treating the dressing of nucleons in finite nuclei in determining the binding energy per particle. It is therefore reasonable to conclude that a careful study of short-range correlations including the full fragmentation of the single-particle strength is relevant for the calculation of the energy per particle in finite nuclei. This has the additional advantage that agreement with data from the (e,e′p) reaction can be used to gauge the quality of the theoretical description in determining the energy per particle. This argument can be turned inside out by noting that an exact representation of the spectroscopic strength must lead to the correct energy per particle according to Eq. (10) in the case of the dominance of two-body interactions.

Pursuing this argument in the case of nuclear matter, while recalling that the empirical saturation density is apparently dominated by SRC, we have calculated nuclear saturation properties focusing solely on the contribution of SRC. The experimental results discussed above demand furthermore that the dressing of nucleons in nuclear matter is then taken into account in order to be consistent with the extensive collection of data from the (e,e′p) reaction that have become available in recent years. The self-consistent calculation of nucleon spectral functions obtained from the contribution to the nucleon self-energy of ladder diagrams which include the propagation of these dressed particles, fulfills this requirement.

It is straightforward to write down the equation that involves the calculation of the effective interaction in nuclear matter obtained from the sum of all ladder diagrams while propagating fully dressed particles. This result is given in a partial wave representation by the following equation

$$\begin{aligned} \langle k|\, \Gamma^{JST}_{LL'}(K,\Omega)\, |k'\rangle &= \langle k|\, V^{JST}_{LL'}(K,\Omega)\, |k'\rangle \\ +\sum_{L''}\int_0^\infty dq\; q^2\; &\langle k|\, V^{JST}_{LL''}(K,\Omega)\, |q\rangle\, g_f^{II}(q;K,\Omega)\, \langle q|\, \Gamma^{JST}_{LL''}(K,\Omega)\, |k'\rangle\,, \end{aligned} \tag{11}$$

where k, k', and q denote relative and K the total momentum involved in the interaction process. Discrete quantum numbers correspond to total spin, S, orbital angular momentum, L, L', L'', and the conserved total angular momentum and isospin, J and T, respectively. The energy Ω and the total momentum K are conserved and act as parameters that characterize the effective two-body interaction in the medium. The critical ingredient in Eq. (11) is the noninteracting propagator g_f^{II} which describes the propagation of the particles in the medium from interaction to interaction. For fully dressed particles this propagator is given by

$$\begin{aligned} g_f^{II}(k_1,k_2;\Omega) = &\int_{\epsilon_F}^\infty d\omega_1 \int_{\epsilon_F}^\infty d\omega_2\; \frac{S_p(k_1,\omega_1)S_p(k_2,\omega_2)}{\Omega-\omega_1-\omega_2+i\eta} \\ &- \int_{-\infty}^{\epsilon_F} d\omega_1 \int_{-\infty}^{\epsilon_F} d\omega_2\; \frac{S_h(k_1,\omega_1)S_h(k_2,\omega_2)}{\Omega-\omega_1-\omega_2-i\eta}, \end{aligned} \tag{12}$$

where individual momenta k_1 and k_2 have been used instead of total and relative momenta as in Eq. (11). The dressing of the particles is expressed by the use of particle and hole spectral functions, S_p and S_h, respectively. The particle spectral function, S_p, is defined as a particle addition probability density in a similar way as the hole spectral function in Eq. (3) for removal. These spectral functions take into account that the particles propagate with respect to the correlated ground state incorporating the presence of high-momentum components in the ground state. This treatment therefore provides the correlated version of the Pauli principle and leads to substantial modification with respect to the Pauli principle effects related to the free Fermi gas. The corresponding propagator is obtained from Eq. (12) by replacing the spectral functions by strength distributions characterized by δ-functions as follows

$$\begin{aligned} S_p(k,\omega) &= \theta(k-k_F)\delta(\omega-\epsilon(k)) \\ S_h(k,\omega) &= \theta(k_F-k)\delta(\omega-\epsilon(k)), \end{aligned} \tag{13}$$

which leads to the Galitski-Feynman propagator including hole-hole as well as particle-particle propagation of particles characterized by single-particle energies $\epsilon(k)$. Discarding the hole-hole propagation then yields the Brueckner ladder diagrams with the usual Pauli operator for the free Fermi gas. The effective interaction obtained by solving Eq. (11) using dressed propagators can be used to construct the self-energy of the particle. With this self-energy the Dyson equation can be solved to generate a new incarnation of the dressed propagator. The process can then be continued by constructing anew the dressed but noninteracting two-particle propagator according to Eq. (12). At this stage one can return to the ladder equation and so on until self-consistency is achieved for the complete Green's function which is then legitimately called a self-consistent Green's function.

While this scheme is easy to present in equations and words, it is quite another matter to implement it. The recent accomplishment of implementing this self-consistency scheme [20] builds upon earlier approximate implementations. The first nuclear-matter spectral functions were obtained for a semirealistic interaction by employing mean-field propagators in the ladder equation [43]. Spectral functions for the Reid interaction were obtained by still employing mean-field propagators in the ladder equation but with the introduction of a self-consistent gap in the single-particle spectrum to take into account the pairing instabilities obtained for a realistic interaction [44, 45]. The first solution of the effective interaction using dressed propagators was obtained by employing a parametrization of the spectral functions [46]. The calculations employing dressed propagators in determining the effective interaction demonstrate that at normal density one no longer runs into pairing instabilities on account of the reduced density of states associated with the reduction of the strength of the quasiparticle pole, z_{k_F}, from 1 in the Fermi gas to 0.7 in the case of dressed propagators. For two-particle propagation this leads to a reduction factor of $z_{k_F}^2$ corresponding to about 0.5 that is strong enough to push even the pairing instability in the $^3S_1$-$^3D_1$ channel to lower densities [47].

The consequences for the scattering process of interacting particles in nuclear matter characterized by phase shifts and cross sections are also substantial and lead to a reduction of the cross section in a wide range of energies [47]. The current implementation of the self-consistent scheme

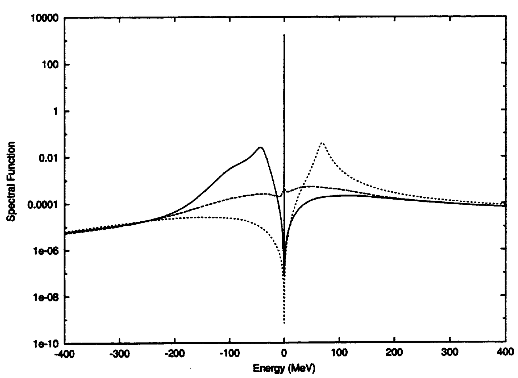

FIG. 3: Self-consistent spectral functions at $k_F = 1.36$ fm^{-1}. Single-particle momenta corresponding to $k = 0$ (solid), k_F (dashed), and 2.1 fm^{-1} (dotted) are shown.

for the propagator across the summation of all ladder diagrams includes a parametrization of the imaginary part of the nucleon self-energy. Employing a representation in terms of two gaussians above and two below the Fermi energy, it is possible to accurately represent the nucleon self-energy as generated by the contribution of relative S-waves (and including the tensor coupling to the $^3D_1$ channel) [20]. Self-consistency at a density corresponding to $k_F = 1.36$ fm^{-1} is achieved in about ten iteration steps, each involving a considerable amount of computer time [20]. It is important to reiterate that this scheme isolates the contribution of short-range correlations to the energy per particle which is obtained from Eq. (10). An important result pertaining to this “second generation” spectral functions is shown in Fig. 3 related to the emergence of a common tail at large negative energy for different momenta. Such a common tail was previously obtained at high energy [41] as a signature of SRC. This common tail may play a significant role in generating some additional binding energy at lower densities. At present, results for two densities corresponding to $k_F = 1.36$ and 1.45 fm^{-1} have been obtained. Self-consistency is achieved for the contribution of the $^1S_0$ and $^3S_1$-$^3D_1$ channels to the self-energy. The other partial wave contributions have been

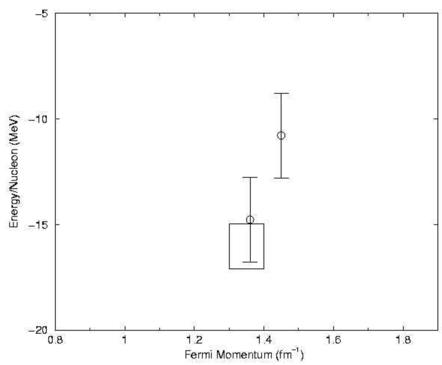

FIG. 4: The energy per particle calculated at two densities. The saturation density for this self-consistent Green's function calculation with the Reid potential is possibly in agreement with the empirical result.

added separately. The corresponding results for the binding energy have been obtained by averaging the parametrizations of the corresponding self-energies. The difference between these results also provides us with a conservative estimate of the lack of self-consistency including higher partial waves. This error estimate is included in Fig. 4 for the energy per particle calculated from the energy (Koltun) sum rule in Eq. (10). These results suggest that it is possible to obtain reasonable saturation properties for nuclear matter provided one only includes SRC in the determination of the equation of state. Clearly, this assertion implies that long-range contributions to the energy per particle need not be considered in explaining nuclear saturation properties. Considerations relevant to this issue are presented elsewhere [48].

5. CONCLUSIONS

A review of experimental data that exhibit clear evidence for the notion that nucleons in nuclei are dressed particles is given. A recent doubt concerning the validity of the interpretation of the spectroscopic strength has been resolved by showing that the same spectroscopic factors can be used to explain data from the (e,e′p) reaction at different values of Q^2. Based on these considerations and the success of theoretical calculations to account for the qualitative features of the single-particle strength distributions, it is suggested that the dressing of nucleons must be taken into account in calculations of the energy per particle. By identifying the dominant contribution of SRC to the empirical saturation density, it is argued that these correlations need to be isolated in the study of nuclear matter. A scheme which fulfills this requirement and includes the propagation of dressed particles, as required by experiment, is outlined. Successful implementation of this scheme has recently been demonstrated [20] for the continuum version. A discrete version has been implemented by the Gent group [49, 50]. These new calculations may lead to new insight into the long-standing problem of nuclear saturation.

ACKNOWLEDGMENTS

This work was supported by the U. S. National Science Foundation under Grant No. PHY-9900713.

REFERENCES

[1] A. E. L. Dieperink and P. K. A. de Witt Huberts, Ann. Rev. Nucl. Part. Sci. **40**, 239 (1990).
[2] I. Sick and P. K. A. de Witt Huberts, Comm. Nucl. Part. Phys. **20**, 177 (1991).
[3] L. Lapikás, Nucl. Phys. **A553**, 297c (1993).
[4] V. R. Pandharipande, I. Sick, and P. K. A de Witt Huberts, Rev. Mod. Phys. **69**, 981 (1997).
[5] W. H. Dickhoff, Phys. Rep. **242**, 119 (1994).
[6] W. H. Dickhoff, in *Nuclear Methods and the Nuclear Equation of State*, ed. M. Baldo (World Scientific, Singapore, 1999) p. 326.
[7] I. Bobeldijk *e*t al., Phys. Rev. Lett. **73**, 2684 (1994).

[8] K. I. Blomqvist *et al.*, Phys. Lett. **B344**, 85 (1985).
[9] H. Müther and W. H. Dickhoff, Phys. Rev. C **49**, R17 (1994).
[10] L. Lapikás, G. van der Steenhoven, L. Frankfurt, M. Strikman, and M. Zhalov, Phys. Rev. C **61** (2000) 064325
[11] A. Bianconi and M. Radici, Phys. Lett. **B363**, 24 (1995).
[12] A. Bianconi and M. Radici, Phys. Rev. C **53**, R563 (1996).
[13] A. Bianconi and M. Radici, Phys. Rev. C **54**, 3117 (1996).
[14] A. Bianconi and M. Radici, Phys. Rev. C **56**, 1002 (1997).
[15] Ye. S Golubeva, L. A. Kondratyuk, A. Bianconi, S. Boffi, and M. Radici, Phys. Rev. C **57**, 2618 (1998).
[16] J. Gao *et al.*, Phys. Rev. Lett. **84**, 3265 (2000).
[17] H. Müther, A. Polls, and W. H. Dickhoff, Phys. Rev. C **51**, 3040 (1995).
[18] A. Polls, M. Radici, S. Boffi, W. H. Dickhoff, and H. Müther, Phys. Rev. C **55**, 810 (1997).
[19] M. Leuschner *et al.*, Phys. Rev. C **49**, 955 (1994).
[20] E. P. Roth, Ph. D. Thesis, Washington University, St. Louis, 2000.
[21] V.M. Galitski and A.B. Migdal, *Sov. Phys. JETP* **34**, 96 (1958).
[22] D. S. Koltun, Phys. Rev. C **9**, 484 (1974).
[23] J. Mougey, Nucl. Phys. **A262**, 461 (1976).
[24] G.J. Wagner, *AIP Conf. Proc.* **142**, 220 (1986).
[25] P. Grabmayr *et al.*, Phys. Lett. **B164**, 15 (1985).
[26] P. Grabmayr, Prog. Part. Nucl. Phys. **29**, 251 (1992).
[27] L. Lapikás, private communication (2000).
[28] M. F. van Batenburg, Ph. D. Thesis, University of Utrecht (2001).
[29] M. G. E. Brand, G. A. Rijsdijk, F. A. Muller, K. Allaart, and W. H. Dickhoff, Nucl. Phys. **A531**, 253 (1991).
[30] G. A. Rijsdijk, K. Allaart, and W. H. Dickhoff, Nucl. Phys. **A550**, 159 (1992).
[31] W. J. W. Geurts, K. Allaart, W. H. Dickhoff, and H. Müther, Phys. Rev. C **53**, 2207 (1996).
[32] C. Barbieri and W. H. Dickhoff, Phys. Rev. C **63**, 034313 (2001).
[33] C. Barbieri and W. H. Dickhoff, contribution to this workshop.
[34] D. Rohe, contribution to this workshop.
[35] M. Zhalov, private communication (2001).
[36] S. Boffi, C. Giusti, F.D. Pacati and M. Radici, *Electromagnetic Response of Atomic Nuclei* (Oxford University Press, Oxford, 1996) vol. 20.
[37] R.J. Glauber, in *Lectures in Theoretical Physics*, eds. W. Brittain and L.G. Dunham (Interscience Publ., New York, 1959) vol. 1.
[38] W. H. Dickhoff, Phys. Rev. C **58**, 2807 (1998).
[39] M. Radici, S. Boffi, S. C. Pieper, and V. R. Pandharipande, Phys. Rev. C **50**, 3010 (1994).
[40] W. H. Dickhoff and H. Müther, Rep. Prog. Phys. **55**, 1947 (1992).
[41] B. E. Vonderfecht, W. H. Dickhoff, A. Polls, and A. Ramos, Phys. Rev. C **44**, R1265 (1991).
[42] B. Frois *et al.*, Phys. Rev. Lett. **38**, 152 (1977).
[43] A. Ramos, Ph.D. Thesis, University of Barcelona, 1988.
[44] B. E. Vonderfecht, Ph.D. Thesis, Washington University, St. Louis, 1991.
[45] B. E. Vonderfecht, W. H. Dickhoff, A. Polls, and A. Ramos, Nucl. Phys. **A555**, 1 (1993).
[46] C. C. Gearhart, Ph.D. Thesis, Washington University, St. Louis, 1994.
[47] W. H. Dickhoff, C. C. Gearhart, E. P. Roth, A. Polls, and A. Ramos, Phys. Rev. C **60**, 064319 (1999).
[48] W. H. Dickhoff, in *Recent Progress in Many-Body Theories*, R. F. Bishop *et al.*, eds. (World Scientific, Singapore) in press.
[49] Y. Dewulf, Ph.D. Thesis, University of Gent (2000).
[50] Y. Dewulf, D. Van Neck, and M. Waroquier, in preparation.